Electron and electron-hole quasiparticle states in a driven quantum contact

Mihajlo Vanević,[1] Julien Gabelli,[2] Wolfgang Belzig,[3] and Bertrand Reulet[4]

[1]*Department of Physics, University of Belgrade, Studentski trg 12, 11158 Belgrade, Serbia*
[2]*Laboratoire de Physique des Solides, Univ. Paris-Sud,*
CNRS, UMR 8502, F-91405 Orsay Cedex, France
[3]*Fachbereich Physik, Universität Konstanz, D-78457 Konstanz, Germany*
[4]*Département de physique, Université de Sherbrooke, Sherbrooke, Québec, J1K 2R1, Canada*
(Dated: January 12, 2016)

We study the many-body electronic state created by a time-dependent drive of a mesoscopic contact. The many-body state is expressed manifestly in terms of single-electron and electron-hole quasiparticle excitations with the amplitudes and probabilities of creation which depend on the details of the applied voltage. We experimentally probe the time dependence of the constituent electronic states by using an analog of the optical Hong-Ou-Mandel correlation experiment where electrons emitted from the terminals with a relative time delay collide at the contact. The electron wave packet overlap is directly related to the current noise power in the contact. We have confirmed the time dependence of the electronic states predicted theoretically by measurements of the current noise power in a tunnel junction under harmonic excitation.

Recent years have seen a tremendous experimental and theoretical progress in the emerging field of electron quantum optics.[1] Following the example of optics, the quantum nature of electronic transport has been demonstrated in electronic Mach-Zehnder interferometer[2] and Hanbury Brown-Twiss[3–5] and Hong-Ou-Mandel[6–8] intensity correlation experiments. Although quantum optics with electrons is in general analogous to the one with photons, there are important distinctions between the two due to differences in particle statistics, vacuum state (Fermi sea vs. photonic vacuum), interaction between electrons, decoherence, etc. In particular, a simple constant voltage source can act as a single-electron turnstile[9] due to the Fermi statistics which is responsible for regular emission of electrons on a time scale h/eV, where e is the electron charge, h is the Planck constant, and V is the dc voltage drop over the conductor.

A step forward towards electron quantum optics has been made recently with the realization of *on-demand* electron sources[4,5,8,10–14] which can create single- to few-particle excitations.[15–21] This facilitates the full control of the quantum state of electrons in mesoscopic conductors and the dynamical control of elementary excitations using suitably tailored voltage pulses.[22] In particular, time-dependent drive creates quasiparticle excitations in the Fermi sea that are single-electron and electron-hole pairs whose number and probability of creation depend on the shape and the amplitude of the applied voltage.[23] Lorentzian pulses $V(t)$ of a quantized area $\int eV(t)dt/h = N$ (N is an integer) are special as they create N electrons above the Fermi level leaving the rest of the Fermi sea unperturbed.[15] Experimentally, the presence of electron-hole pairs in the system can be seen in the zero-frequency photon-assisted current noise power which is increased with respect to the dc noise level.[24–26] More recently, quantum noise oscillations have been observed in a driven tunnel junction[27] and noise

spectroscopy using a more complex biharmonic voltage drive has been carried out approximating Lorentzian pulses.[14] The creation of single-electron excitations has been realized experimentally[12] and the resulting quantum states have been reconstructed using the quantum state tomography.[13]

Even though the progress has been exceptional, a fundamental question remains: What is the many-body electronic state created in the Fermi sea by a voltage drive? In this article we find that the many-body state is:

$$|\Psi\rangle = \hat{C}^\dagger \prod_k \left(\sqrt{1 - p_k} + i\sqrt{p_k}\hat{A}_k^\dagger\hat{B}_k \right) |F\rangle. \qquad (1)$$

Here, $\hat{C}^\dagger = \int d\mathcal{E}\mathrm{v}^*(\mathcal{E})\hat{c}^\dagger(\mathcal{E})$ is the creation operator of a single-electron quasiparticle state, $\hat{A}_k^\dagger = \int d\mathcal{E}\mathrm{u}_+^{k*}(\mathcal{E})\hat{c}^\dagger(\mathcal{E})$ and $\hat{B}_k = \int d\mathcal{E}\mathrm{u}_-^k(\mathcal{E})\hat{c}(\mathcal{E})$ are the operators that create electron and hole from the electron-hole pair, $\hat{c}^\dagger$ ($\hat{c}$) are the electron creation (annihilation) operators, and $|F\rangle$ is the filled Fermi sea. [We have assumed, for simplicity, that there is one single-electron quasiparticle created per period, which is the case for $eV_{\mathrm{dc}}/\hbar\omega = 1$ where V_{dc} is a dc voltage component and ω is the frequency of the drive.] In addition to a single-electron excitation, there is a number of electron-hole pairs created in the system (labelled by $k = 1, 2, \ldots$) due to the ac voltage component. The probabilities of the electron-hole pair creations p_k and the single-electron and electron-hole quasiparticle amplitudes $\mathrm{v}(\mathcal{E})$ and $\mathrm{u}_\pm^k(\mathcal{E})$ depend on the properties of the applied voltage. For optimal Lorentzian drive, there are no electron-hole pairs created ($p_k = 0$) and the state has only single-electron excitations, as expected.

Quasiparticle amplitudes v and $\mathrm{u}_\pm^k$ in Eq. (1) also give the time-dependent probabilities of single-electron and electron-hole pair creations, $|\mathrm{v}(t)|^2$ and $|\mathrm{u}_\pm^k(t)|^2$. The time dependence of the wave functions that constitute the many-body state in Eq. (1) can be probed by an elec-

tronic analog of the optical Hong-Ou-Mandel correlation experiment,[28] where electrons emitted from two terminals with a relative time delay collide at the contact.[29,30] When the wave packets arrive at the contact simultaneously, their transmission in the same output channel is blocked by the Pauli principle, which suppresses the current fluctuations. The magnitude of the noise suppression is proportional to the wave-packet overlap at the contact. In the present paper, we have performed this experiment and measured current noise power in a tunnel junction driven by harmonic time-dependent voltage. We have found that the correlation noise as a function of a time delay is in agreement with the theoretically predicted quasiparticle amplitudes of electrons and electron-hole pairs in Eq. (1).

Let us consider a generic quantum contact with spin-degenerate transmission eigenvalues T_n that are independent of energy. The contact is driven by a voltage $V(t) = V_{dc} + V_{ac}(t)$, where V_{dc} is a constant dc offset and $V_{ac}(t)$ is a periodic ac voltage component with zero average and the period $T = 2\pi/\omega$. The cumulant generating function of the charge transfer statistics is given by[15,23] $\mathcal{S}(\chi) = 2\sum_n \text{Tr}\ln[1 + f_L(1 - f_R)T_n(e^{i\chi} - 1) + (1 - f_L)f_R T_n(e^{-i\chi} - 1)]$. Here, $f_{L(R)}$ are generalized electronic distribution functions in the left (right) terminal which depend on two time or energy arguments: $f_L \equiv \tilde{f} = e^{-i\phi(t')}f_{V_{dc}}(t' - t'')e^{i\phi(t'')}$, $f_R \equiv f(t' - t'')$. Here, $\phi(t) = \int_0^t dt' eV_{ac}(t')/\hbar$ (hereafter $\hbar = 1$), $f(\mathcal{E}) = 1 - \theta(\mathcal{E})$, and $f_{V_{dc}}(\mathcal{E}) = f(\mathcal{E} - eV_{dc})$, where θ is the step function. Function $\tilde{f}$ couples only energies which differ by an integer multiple of ω, $\tilde{f}(\mathcal{E}', \mathcal{E}'') = \sum_k a_k a_{k+m-n}^* f(\mathcal{E}' - k\omega - eV_{dc})$, where $\mathcal{E}' = \epsilon + n\omega$, $\mathcal{E}'' = \epsilon + m\omega$ ($0 < \epsilon < \omega$), and $a_n = (1/T)\int_0^T dt e^{-i\phi(t)}e^{in\omega t}$. In particular, diagonal components of $\tilde{f}$ are given by $\tilde{f}(\mathcal{E}) = \sum_{n=-\infty}^\infty |a_n|^2 f(\mathcal{E} - n\omega - eV_{dc})$. Function $\tilde{f}(\mathcal{E}', \mathcal{E}'')$ is the generalized distribution function of a driven quantum contact in energy representation, while $\tilde{f}(\mathcal{E})$ is the stationary nonequilibrium electronic distribution which is realized in the junction due to time-dependent drive.

Next we obtain a decomposition of $\tilde{f}$ into single-electron and electron-hole states. As shown in Ref. 23, the notion of single-electron and electron-hole pair excitations is related to the eigenproblem of $\{h, \tilde{h}\} \equiv h\tilde{h} + \tilde{h}h$, where $f = (1 - h)/2$ and $\tilde{f} = (1 - \tilde{h})/2$. For integer $eV_{dc}/\omega = N$, there is an N-dimensional subspace of $\{h, \tilde{h}\}$ spanned by N *special* vectors that are eigenvectors of *both* h and $\tilde{h}$: $\tilde{h}\mathbf{v} = -\mathbf{v}$, $h\mathbf{v} = \mathbf{v}$. This subspace corresponds to N *electrons* injected to the contact per voltage cycle. In addition, the operator $\{h, \tilde{h}\}$ has a series of two-dimensional subspaces that are spanned by vectors $\mathbf{v}_\alpha$ and $\mathbf{v}_{-\alpha} \equiv h\mathbf{v}_\alpha$ which are given by $h\tilde{h}\mathbf{v}_\alpha = e^{i\alpha}\mathbf{v}_\alpha$. The spaces $\text{span}(\mathbf{v}_{\alpha_k}, \mathbf{v}_{-\alpha_k})$ correspond to the *electron-hole pairs* (labelled by $k = 1, 2, \ldots$) created per voltage cycle with the probabilities $p_k = \sin^2(\alpha_k/2)$.

At zero temperature, f and $\tilde{f}$ commute with $\{h, \tilde{h}\}$, that is, they reduce in the single-electron and electron-hole pair subspaces of $\{h, \tilde{h}\}$. For simplicity, we restrict

our consideration to the case $N = 1$ in which there is only one electron injected per voltage cycle. The eigenproblem of $\{h, \tilde{h}\}$ defines a resolution of the identity $|\mathbf{v}\rangle\langle\mathbf{v}| + \sum_k \hat{P}_k = 1$, where $|\mathbf{v}\rangle\langle\mathbf{v}|$ is the projector on the single-electron state and $\hat{P}_k = |\mathbf{v}_{\alpha_k}\rangle\langle\mathbf{v}_{\alpha_k}| + |\mathbf{v}_{-\alpha_k}\rangle\langle\mathbf{v}_{-\alpha_k}|$ are the projectors on the electron-hole subspaces. This defines a decomposition of $\tilde{f}$ into single-electron and electron-hole contributions,

$$\tilde{f} = |\mathbf{v}\rangle\langle\mathbf{v}| + \sum_k \tilde{f}_k, \qquad (2)$$

where $\tilde{f}_k = \tilde{f}\hat{P}_k$. Similarly, for the completely filled Fermi sea of the right lead we have $f = \sum_k f_k$ where $f_k = f\hat{P}_k$. The single-electron state $\mathbf{v}$ is given by $\tilde{f}\mathbf{v} = \mathbf{v}$ and $f\mathbf{v} = 0$.

The first term $|\mathbf{v}\rangle\langle\mathbf{v}|$ in Eq. (2) describes a single-electron state injected to the contact while $\tilde{f}_k$ describe the electron-hole pairs. By taking diagonal in time components of $|\mathbf{v}\rangle\langle\mathbf{v}|$ or $\tilde{f}_k$ we gain information on the time dependence of single-electron and electron-hole pair wave functions. Similarly, by taking diagonal in energy components we gain information on the single-electron and electron-hole pair contributions in the overall distribution function. Indeed, for Lorentzian pulses $V_{Lor}(t)$ that carry a single charge quantum per cycle, there are no additional electron-hole excitations and the time-dependent probability of the single-electron injection is proportional to the drive, $|\mathbf{v}(t)|^2 = eV_{Lor}(t)/\omega$. This is no longer true for a general time-dependent drive where $|\mathbf{v}(t)|^2 \neq eV(t)/\omega$ due to the presence of electron-hole pairs.

Before we proceed with the specific examples, let us bring $\tilde{f}_k$ in a more transparent form. States $\mathbf{v}_\alpha$ and $\mathbf{v}_{-\alpha}$ in general possess both positive and negative energy components. We can make a rotation of the basis $\mathbf{u}_\pm = (\mathbf{v}_\alpha \pm \mathbf{v}_{-\alpha})/\sqrt{2}$ in the subspace $\text{span}(\mathbf{v}_\alpha, \mathbf{v}_{-\alpha})$ such that new basis vectors $\mathbf{u}_+$ and $\mathbf{u}_-$ possess non-zero components only for $\mathcal{E} > 0$ and $\mathcal{E} < 0$, respectively. Using $h\mathbf{v}_\alpha = \mathbf{v}_{-\alpha}$ it is straightforward to check that $f\mathbf{u}_+ = 0$ and $f\mathbf{u}_- = \mathbf{u}_-$, as required. The projector $\hat{P}_k$ in the new basis reads $\hat{P}_k = |\mathbf{u}_+^k\rangle\langle\mathbf{u}_+^k| + |\mathbf{u}_-^k\rangle\langle\mathbf{u}_-^k|$. We obtain

$$\tilde{f}_k = p_k|\mathbf{u}_+^k\rangle\langle\mathbf{u}_+^k| + q_k|\mathbf{u}_-^k\rangle\langle\mathbf{u}_-^k| - i\sqrt{p_k q_k}|\mathbf{u}_+^k\rangle\langle\mathbf{u}_-^k| + \text{h.c.} \qquad (3)$$

and $f_k = |\mathbf{u}_-^k\rangle\langle\mathbf{u}_-^k|$ ($q_k = 1 - p_k$). Decomposition of $\tilde{f}$ in Eqs. (2) and (3) enables us to identify the many-body electronic state in the left lead created by the drive. We find that $\tilde{f}(\mathcal{E}', \mathcal{E}'') = \langle\Psi|\hat{c}^\dagger(\mathcal{E}')\hat{c}(\mathcal{E}'')|\Psi\rangle$ where $|\Psi\rangle$ is given in Eq. (1). The right lead is assumed to be unperturbed and serves as a reference.

The physical meaning of the amplitudes $\mathbf{u}_\pm$ is manifest in Eq. (1) and can further be elaborated by taking diagonal in energy components of $\tilde{f}_k$. For $\mathcal{E} > 0$ we find that $\tilde{f}_k(\mathcal{E}) = p_k|\mathbf{u}_+^k(\mathcal{E})|^2$. For $\mathcal{E} < 0$ it is more convenient to consider the distribution of holes $\tilde{f}_k^{(h)} \equiv \hat{P}_k - \tilde{f}_k$: $\tilde{f}_k^{(h)}(\mathcal{E}) = p_k|\mathbf{u}_-^k(\mathcal{E})|^2$. Therefore, $\mathbf{u}_+$ and $\mathbf{u}_-$ describe electrons and holes generated in the system. This notion is also supported in time domain with diagonal in

time components $\tilde{f}_k(t) = |\mathbf{u}_+^k(t)|^2$ and $\tilde{f}_k^{(h)}(t) = |\mathbf{u}_-^k(t)|^2$. Apart from electron and hole states on the diagonal, $\tilde{f}_k$ contains also the off-diagonal terms proportional to $\sqrt{p_k(1-p_k)}$ that are responsible for mixing of the two, see Eq. (3). Electron-hole pairs with $p_k \approx 0$ give no contribution to the transport. On the other hand, for $p_k \approx 1$ the electron and the hole from a pair are practically decoupled from each other (off-diagonal mixing terms vanish), cf. Eq. (1). This can also be seen in the cumulant generating function which becomes a sum of independent electron and hole contributions.[23]

Next we study single-electron and electron-hole pair states for different voltage drives, see Fig. 1. Let us consider a harmonic drive $V(t) = V_{dc} + V_0\cos(\omega t)$ where the dc offset $N = 1$ is kept fixed while the amplitude V_0 of the ac component is varied. The drive is characterized by the coefficients $a_n = J_n(eV_0/\omega)$ where J_n are the Bessel functions of the first kind. In addition to a single-electron state, there are also electron-hole pairs created in the system and they become more relevant for transport as the amplitude V_0 is increased. For $eV_0/\omega \lesssim 2$ there is only one electron-hole pair with $p_1 = (1/2)(\sum_{n=-\infty}^{\infty} |n + N| |a_n|^2 - N)$ in addition to the single-electron state injected. For $eV_0/\omega \ll 1$, the probability p_1 practically vanishes and only a single-electron state remains. Time dependence of the single-electron wave packet $|\mathbf{v}(t)|^2$ is shown in Figs. 1(a–c) in comparison to the voltage drive $V(t)$. We find that for ac drive amplitudes much smaller than dc voltage bias, the electron-hole pair creation is not effective and the single-electron wave packet coincides with the drive, $|\mathbf{v}(t)|^2 \approx eV(t)/\omega$, see Fig. 1(a). For ac drive amplitudes comparable to or larger than dc offset, the single-electron wave packet differs from the drive, see Figs. 1(b,c). In that case the electron-hole pairs become important and their wave packets together with the single-electron one ensure $I(t) = GV(t)$, where $G = (e^2/\pi)\sum_n T_n$. The wave functions $|\mathbf{u}_\pm(t)|^2$ of an electron-hole pair are shown in Fig. 1(d) for the drive without dc bias ($N = 0$) which does not create single-electron states. The nonequilibrium distribution functions $\tilde{f}(\mathcal{E})$ for the voltage drives at hand are shown in Figs. 1(e–h) (solid curves) together with the approximations (dash-dotted curves) computed using the most important single-electron or electron-hole pair states in Figs. 1(a–d). From Figs. 1(b,c,f,g) we find that electron-hole pairs give a more significant contribution in time domain than in $\tilde{f}(\mathcal{E})$. Indeed, while the single-electron wave packets clearly differ from the voltage drive, the distribution function is nevertheless to a good accuracy given by a single-electron state, $\tilde{f}(\mathcal{E}) \approx |\mathbf{v}(\mathcal{E})|^2 + \theta(-\mathcal{E})$. This is because for small p_k, the electron-hole pair functions $\tilde{f}_k$ in Eq. (3) have dominant off-diagonal electron-hole mixing components (proportional to $\sqrt{p_k}$) which do not contribute to diagonal in energy distribution $\tilde{f}_k(\mathcal{E})$, while they do contribute to diagonal in time distribution.

The time dependence of the electronic wave functions can be accessed experimentally by studying current noise
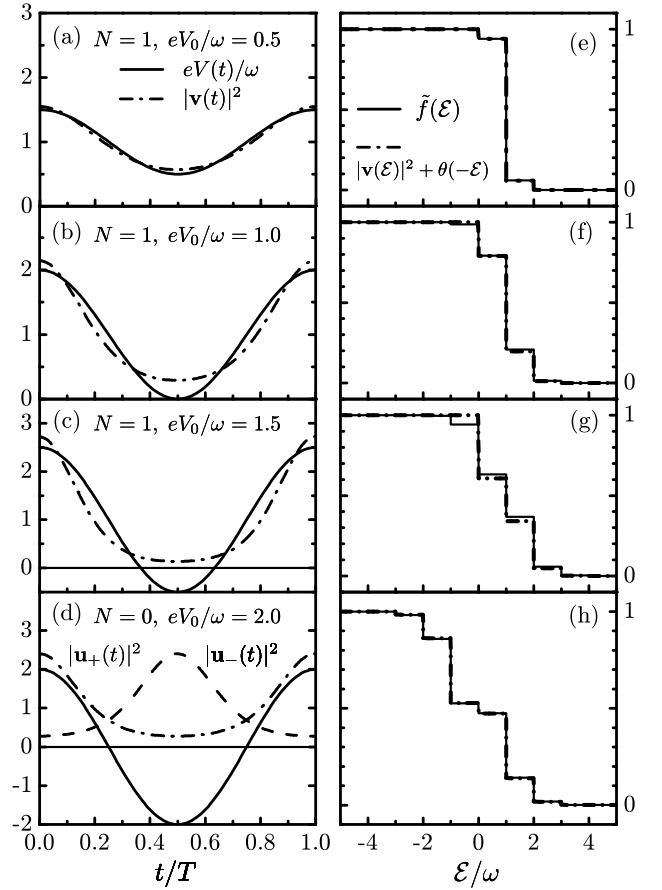

FIG. 1. (a)–(d) Wave functions $|\mathbf{v}(t)|^2$ and $|\mathbf{u}_\pm(t)|^2$ of electrons and electron-hole pairs created by different voltage drives $V(t)$. The corresponding nonequilibrium distribution functions are shown in the panels (e)–(h).

power in a setup where two voltage drives with time shift τ are applied to the terminals, $V_L(t) = V_{dc} + V_0\cos(\omega t)$ and $V_R(t) = V_L(t - \tau)$. This can be viewed as the electronic analog of the optical Hong-Ou-Mandel (HOM) experiment[28] in which electron wave packets emitted from the terminals with time delay τ collide at the contact.[29,30] In the analog HOM experiment, V_{dc} is the static bias voltage between the input and output ports, which is in our case only defined with respect to virtual output terminals. Because of the gauge invariance, our two-terminal setup is formally equivalent to the case of the voltage $\delta V(t) = V_L(t) - V_R(t)$ applied only to the left lead with the right lead unperturbed. The current noise power as a function of the time delay τ then reads $S_2(\tau) = S_0 \sum_{n=-\infty}^{\infty} |n| J_n^2[(2eV_0/\omega)\sin(\omega\tau/2)]$ where $S_0 = (e^2\omega/\pi)\sum_n T_n R_n$ with $R_n = 1 - T_n$.[31] To express $S_2(\tau)$ in terms of the overlap of the wave packets, we proceed as follows. From the cumulant generating function $\mathcal{S}$ we find $S_2(\tau) \propto \text{Tr}[(f_L - f_R)^2] = \text{Tr}[(f_L - f)^2 + (f_R - f)^2 - 2(f_L - f)(f_R - f)]$ where we have introduced the distribution function f of the unperturbed Fermi sea. Here, the first two terms on
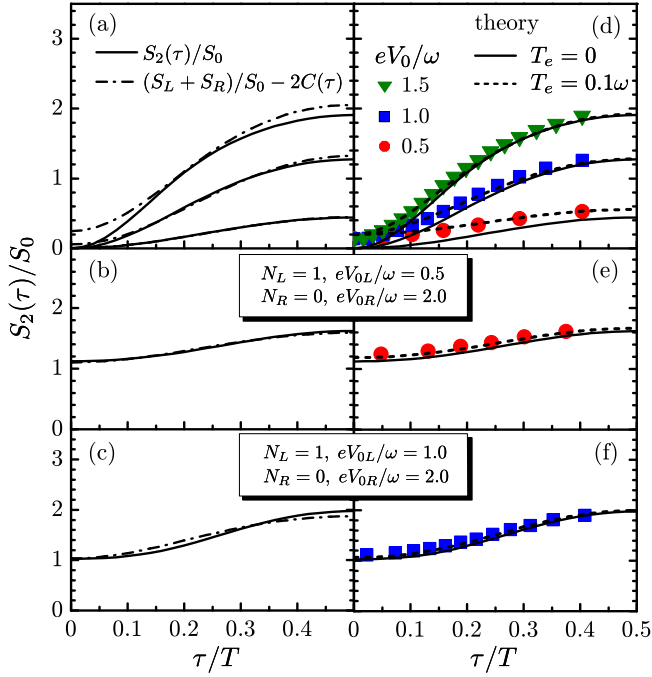
FIG. 2. Noise S_2 as a function of a time delay τ between two harmonic signals applied at the leads: (a) $N = 1$ and $eV_0/\omega = 1.5, 1, 0.5$ (solid lines, top to bottom), (b) left lead: $N = 1$, $eV_0/\omega = 0.5$; right lead: $N = 0$, $eV_0/\omega = 2$, (c) left lead: $N = 1$, $eV_0/\omega = 1$; right lead: $N = 0$, $eV_0/\omega = 2$. Approximations for S_2 calculated using the overlaps $C(\tau)$ of the wave functions depicted in Fig. 1 are shown for comparison (dash-dotted lines). (d)–(f) $S_2(\tau)$ measured in a tunnel junction at $\omega/2\pi = 20\,\mathrm{GHz}$ (symbols). Theoretical results are shown for temperatures $T_e = 0$ (solid lines) and $T_e = 0.1\omega$ (dashed lines).

the right-hand side give the noise when the voltage is applied to one lead only while the other lead is unperturbed. Both terms give the same contribution to the noise $S_L = S_R = S_0 \sum_{n=-\infty}^{\infty} |n + N| J_n^2(eV_0/\omega)$ independent of τ. The term $-\mathrm{Tr}[2(f_L - f)(f_R - f)]$ gives the noise suppression due to wave packet overlap at the contact. The noise reads

$$S_2(\tau)/S_0 = (S_L + S_R)/S_0 - 2C(\tau), \qquad (4)$$

where $C(\tau)$ is the overlap that we compute using Eqs. (2) and (3). For $p_k \ll 1$, the dominant contribution in $C(\tau)$ is the overlap of the single-electron wave functions $\mathbf{v}(t)$ and $\mathbf{v}'(t) \equiv \mathbf{v}(t - \tau)$ injected from the leads,

$$C(\tau) \approx |\langle \mathbf{v} | \mathbf{v}' \rangle|^2 = \left| \int_0^T \frac{dt}{T} \, \mathbf{v}^*(t) \mathbf{v}(t - \tau) \right|^2. \qquad (5)$$

The noise $S_2(\tau)$ is shown in Fig. 2(a) together with the

noise computed using the most important single-electron wave packet overlap in Eq. (5). The corresponding voltage drives and the electron wave functions are shown in Figs. 1(a–c).

So far we have analyzed the single-electron wave packets. To probe the electron-hole states we can use the voltages $V_L(t) = V_{\mathrm{dc}} + V_{0L}\cos(\omega t)$ and $V_R(t) = V_{0R}\cos[\omega(t - \tau)]$ which create single-electron wave packets at the left lead and electron-hole pairs at the right lead. The current noise power in this case is given by $S_2(\tau) = S_0 \sum_{n=-\infty}^{\infty} |n + N| \, J_n^2[eV_0(\tau)/\omega]$, where $V_0(\tau) = [V_{0L}^2 + V_{0R}^2 - 2V_{0L}V_{0R}\cos(\omega\tau)]^{1/2}$. In terms of the overlap, the noise is given by Eq. (4) where $S_L = S_0 \sum_n |n + N| J_n^2(eV_{0L}/\omega)$, $S_R = S_0 \sum_n |n| J_n^2(eV_{0R}/\omega)$, and $C(\tau) \approx p_R |\langle \mathbf{v} | \mathbf{u}'_+ \rangle|^2$ is the overlap between the single-electron state $\mathbf{v}$ at the left lead and the electron part $\mathbf{u}'_+$ of the most dominant electron-hole pair $\mathbf{u}'_\pm$ generated in the right lead ($p_R = 0.630$). The noise $S_2(\tau)$ in this case is shown in Figs. 2(b,c); the corresponding single-electron and electron-hole wave functions are shown in Figs. 1(a,b,d).

To verify this picture, we have measured $S_2(\tau)$ in a tunnel junction under harmonic excitation with $\omega/2\pi = 20\,\mathrm{GHz}$ (see Supplemental material). Experimental results are shown in Figs. 2(d–f), in agreement with the current noise power obtained theoretically. This proves that single-electron and electron-hole excitations in Eq. (1) can be created by time-dependent voltage and accessed experimentally in a noise correlation experiment which measures the overlap of the electronic wave functions.

In conclusion, we have obtained the many-body electronic state created by a time dependent drive of a quantum contact in terms of single-electron and electron-hole quasiparticle excitations. We have confirmed our theoretical predictions by probing the constituent quasiparticle states in a HOM-type experiment on a tunnel junction. The knowledge of the many-body state opens a way of engineering the required time profile or energy distribution of single-electron and electron-hole excitations. Since the electronic state in a conductor determines the electromagnetic field it generates,[32] our work can be used to produce non-classical states of electromagnetic field, such as squeezed or entangled photonic states that have been observed recently.[33,34]

[1] E. Bocquillon, V. Freulon, F. D. Parmentier, J.-M. Berroir, B. Placais, C. Wahl, J. Rech, T. Jonckheere, T. Martin, C. Grenier, D. Ferraro, P. Degiovanni, and G. Feve, Annalen der Physik **526**, 1 (2014).

[2] Y. Ji, Y. Chung, D. Sprinzak, M. Heiblum, D. Mahalu, and H. Shtrikman, Nature **422**, 415 (2003).

[3] M. Henny, S. Oberholzer, C. Strunk, T. Heinzel, K. Ensslin, M. Holland, and C. Schönenberger, Science **284**, 296 (1999).

[4] E. Bocquillon, F. D. Parmentier, C. Grenier, J.-M. Berroir, P. Degiovanni, D. C. Glattli, B. Plaçais, A. Cavanna, Y. Jin, and G. Fève, Phys. Rev. Lett. **108**, 196803 (2012).

[5] N. Ubbelohde, F. Hohls, V. Kashcheyevs, T. Wagner, L. Fricke, B. Kästner, K. Pierz, H. W. Schumacher, and R. J. Haug, Nature Nanotechnology **10**, 46 (2015).

[6] R. C. Liu, B. Odom, Y. Yamamoto, and S. Tarucha, Nature **391**, 263 (1998).

[7] I. Neder, N. Ofek, Y. Chung, M. Heiblum, D. Mahalu, and V. Umansky, Nature **448**, 333 (2007).

[8] E. Bocquillon, V. Freulon, J.-M. Berroir, P. Degiovanni, B. Plaçais, A. Cavanna, Y. Jin, and G. Fève, Science **339**, 1054 (2013).

[9] K. Thibault, J. Gabelli, C. Lupien, and B. Reulet, Phys. Rev. Lett. **114**, 236604 (2015).

[10] J. Gabelli, G. Fève, J.-M. Berroir, B. Plaçais, A. Cavanna, B. Etienne, Y. Jin, and D. C. Glattli, Science **313**, 499 (2006).

[11] G. Fève, A. Mahé, J.-M. Berroir, T. Kontos, B. Plaçais, D. C. Glattli, A. Cavanna, B. Etienne, and Y. Jin, Science **316**, 1169 (2007).

[12] J. Dubois, T. Jullien, F. Portier, P. Roche, A. Cavanna, Y. Jin, W. Wegscheider, P. Roulleau, and D. C. Glattli, Nature **502**, 659 (2013).

[13] T. Jullien, P. Roulleau, B. Roche, A. Cavanna, Y. Jin, and D. C. Glattli, Nature **514**, 603 (2014).

[14] J. Gabelli and B. Reulet, Phys. Rev. B **87**, 075403 (2013).

[15] J. Keeling, I. Klich, and L. S. Levitov, Phys. Rev. Lett. **97**, 116403 (2006); D. A. Ivanov, H. W. Lee, and L. S. Levitov, Phys. Rev. B **56**, 6839 (1997); L. S. Levitov, H. Lee, and G. B. Lesovik, J. Math. Phys. **37**, 4845 (1996).

[16] J. Keeling, A. V. Shytov, and L. S. Levitov, Phys. Rev. Lett. **101**, 196404 (2008)

[17] M. Moskalets, P. Samuelsson, and M. Büttiker, Phys. Rev. Lett. **100**, 086601 (2008).

[18] S. Ol'khovskaya, J. Splettstoesser, M. Moskalets, and M. Büttiker, Phys. Rev. Lett. **101**, 166802 (2008).

[19] F. D. Parmentier, E. Bocquillon, J.-M. Berroir, D. C. Glattli, B. Plaçais, G. Fève, M. Albert, C. Flindt, and M. Büttiker, Phys. Rev. B **85**, 165438 (2012).

[20] G. Haack, M. Moskalets, and M. Büttiker, Phys. Rev. B **87**, 201302(R) (2013).

[21] M. Moskalets, Phys. Rev. B **91**, 195431 (2015).

[22] M. Vanević and W. Belzig, Phys. Rev. B **86**, 241306(R) (2012).

[23] M. Vanević, Y. V. Nazarov, and W. Belzig, Phys. Rev. Lett. **99**, 076601 (2007); Phys. Rev. B **78**, 245308 (2008).

[24] R. J. Schoelkopf, A. A. Kozhevnikov, D. E. Prober, and M. J. Rooks, Phys. Rev. Lett. **80**, 2437 (1998).

[25] A. A. Kozhevnikov, R. J. Schoelkopf, and D. E. Prober, Phys. Rev. Lett. **84**, 3398 (2000).

[26] L.-H. Reydellet, P. Roche, D. C. Glattli, B. Etienne, and Y. Jin, Phys. Rev. Lett. **90**, 176803 (2003).

[27] G. Gasse, L. Spietz, C. Lupien, and B. Reulet, Phys. Rev. B **88**, 241402(R) (2013).

[28] C. K. Hong, Z. Y. Ou, and L. Mandel, Phys. Rev. Lett. **59**, 2044 (1987); Z. Y. Ou and L. Mandel, **62**, 2941 (1989).

[29] J. Dubois, J. Thibaut, C. Grenier, P. Degiovanni, P. Roulleau, and D. C. Glattli, Phys. Rev. B **88**, 085301 (2013).

[30] C. Grenier, R. Herve, E. Bocquillon, F. D. Parmentier, B. Placais, J. M. Berroir, G. Feve, and P. Degiovanni, New Journal of Physics **13**, 093007 (2011).

[31] We note that the noise is independent of V_{dc} in our case. In a true HOM setup the finite dc-bias is applied between input and output leads. This will lead to the transmission of single-electron states and add a constant noise background independent of the time-dependent drive.

[32] F. Qassemi, A. L. Grimsmo, B. Reulet, and A. Blais, arXiv:1507.00322 (2015).

[33] A. Bednorz, C. Bruder, B. Reulet, and W. Belzig, Phys. Rev. Lett. **110**, 250404 (2013); G. Gasse, C. Lupien, and B. Reulet, **111**, 136601 (2013).

[34] J.-C. Forgues, C. Lupien, and B. Reulet, Phys. Rev. Lett. **114**, 130403 (2015).

Supplemental Material for
"Electron and electron-hole quasiparticle states in a driven quantum contact"

I. EXPERIMENTAL SETUP

In order to generate a single electron state by applying a periodic voltage bias $[V(t) = V_{dc} + V_0 \cos(\omega t)]$ on a quantum contact, the experiment has to be performed in the quantum regime, that is, $\hbar\omega \gg k_B T$ at low temperature $T \sim 100\,\mathrm{mK}$ and high frequency $\omega/2\pi = 20\,\mathrm{GHz}$. To provide a good matching to the coaxial cable and avoid reflection of the ac excitation, one use a $R_0 \simeq 50\Omega$ Al/Al oxide/Al tunnel junction as a quantum contact. Indeed, the time dependence of single-electron and electron-hole wave functions in a driven quantum contact does not depend either on the number nor on the transmission of its conduction channels. The experiment is performed in a dilution refrigerator where a $0.1\,\mathrm{T}$ perpendicular magnetic field is applied to turn the Al normal. A bias tee, sketched in Fig. 3 by an inductor and a capacitor, allows to separate the dc bias voltage applied on the junction from the high frequency part of the setup. The ac excitation is imposed through a directional coupler and the current fluctuations emerging from the sample are amplified by a low noise cryogenic amplifier (noise temperature $T_N \sim 7\,K$). Current fluctuations are band-pass filtered between $\Delta f = 1.2 - 3\,\mathrm{GHz}$ and the noise power density S_2 integrated over the bandwidth of the filter is measured with a broadband square law detectors as a function of dc bias voltage for different ac excitations. The derivative of the noise $\partial S_2/\partial(eV_{dc})$ is measured with an additional $77\,\mathrm{Hz}$, small voltage modulation and a usual lock-in detection. The base temperature T_{ph} of the dilution refrigerator significantly increases with increasing ac voltage. It implies an increase of the electronic temperature T_e which is measured by fitting the data of $\partial S_2/\partial(eV_{dc})$ (see inset in Fig. 3).

II. CALIBRATION PROCEDURE

The current noise power associated to the collision of two single-electron wave packets generated by two voltage drives $V_L(t) = V_{dc} + V_0\cos(\omega t)$ and $V_R(t) = V_L(t-\tau)$ is simply given by the harmonic photon-assisted noise with $V_{ac} = 2eV_0 \sin(\omega\tau/2)$:
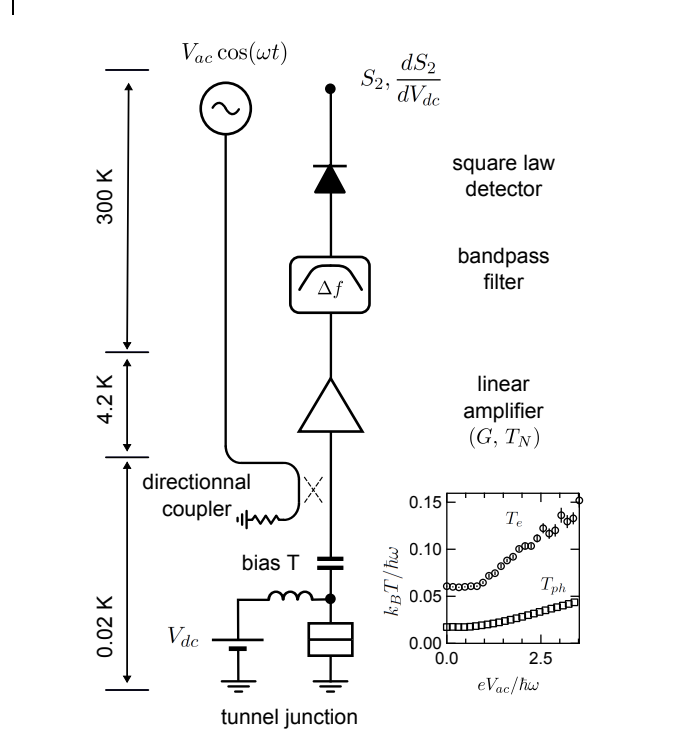
FIG. 3. Experimental setup for the measurement of the current noise power in a tunnel junction under harmonic excitation.

$$S_2(eV_{dc}, eV_{ac}) = \sum_{n=-\infty}^{+\infty} J_n^2\left(\frac{eV_{ac}}{\hbar\omega}\right) \frac{R_0(eV_{dc} + n\hbar\omega)}{\mathrm{th}\left(\frac{eV_{dc}+n\hbar\omega}{2k_B T_e}\right)},$$

where $S_0 = \hbar\omega/R_0$ and T_e is the temperature of electrons. Since we measure the amplified current noise power density $G(S_2 + 2k_B T_N)\Delta f$ and we want to extract its absolute value S_2, we need to calibrate the total gain G of the measurement setup (including the amplification and the cable attenuation), the noise added by the amplifier T_N and the coupling $eV_{ac}/\hbar\omega$ between the sample and the excitation line. This calibration has been done in two stages: (i) $G\Delta f$ and T_N are deduced from the high voltage dependence of the noise power which is simply given by the classical shot noise: